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One-loop calculations of hyperon polarizabilities under the large N_c consistency condition

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Abstract

The spin-averaged electromagnetic polarizabilities of the hyperons Λ and Σ are calculated within the one-loop approximation by use of the dispersion theory. The photon and meson couplings to hyperons are determined so as to satisfy the large N_c consistency condition. It is shown that in order for the large N_c consistency condition to hold exotic hyperon states such as Σ^{**} with $I = 2$ and $J = 3/2$ are required in the calculation of the magnetic polarizability of the Σ state.

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I. INTRODUCTION

Beyond the spin averaged electromagnetic polarizabilities of the nucleon, the spin polarizabilities [1] have recently attracted theoretical attention, because these quantities serve as a crucial test of the low energy effective theories. Using the heavy baryon chiral perturbation theories (HBChPT) [2–4], the spin polarizabilities have been calculated and compared with the multipole analyses [5–7]. The spin polarizabilities are also calculated by using the dispersion theory, where the imaginary parts are given by the Born terms of the one pion photoproduction amplitude [8,9]. The dispersion theory with the Born terms is a method to calculate loop diagrams [10–12], and it reproduces almost the same results by HBChPT up to $\mathrm{O}(p^3)$ or $\mathrm{O}(\varepsilon^3)$, but it includes partially higher chiral order diagrams than $\mathrm{O}(p^3)$; for example Ref. [9] gives the forward spin polarizability $\gamma_0 = -0.4 \times 10^{-4}\,\mathrm{fm}^4$, while HBChPT up to $\mathrm{O}(\epsilon^3)$ does $\gamma_0 = 2.0 \times 10^{-4}\,\mathrm{fm}^4$ [4].

As to hyperons the spin-averaged polarizabilities have so far been studied in the quark model [13], the SU(3) extension of the HBChPT results [14] and the bound-state soliton model [15], but the study of hyperon polarizabilities is quite insufficient, because the hyperon polarizabilities involve much more physical contents than that of the nucleon. Further, since measurements of $\gamma\Sigma$ interactions are planned [16], detailed and comprehensive studies will be required.

In this paper we calculate the spin-averaged polarizabilities of the Λ and Σ hyperons within the one-loop approximation by applying the dispersion theory to the Compton scattering amplitude, where the imaginary parts of the amplitudes are given by the Born terms of the pion and kaon photoproduction amplitudes. The coupling constants of the photon and meson to the nucleon and hyperons in the meson photoproduction amplitudes are given by the spin-flavor symmetry which leads to the large N_c consistency condition. We refer sometimes to the bound state approach to strangeness in the chiral soliton model [17,18], because it is an explicit model realizing the spin-flavor symmetry for baryons in large N_c QCD [19,20], and the results are shared with any large N_c baryon theories at leading order in the $1/N_c$ expansion. We call the model as the bound kaon-soliton model(BKSM) hereafter.

The Born terms in the pion and kaon photoproduction amplitudes with the electric coupling of the photon, which we call the electric Born terms, contribute to the electric and magnetic polarizabilities. The polarizabilities by the pion electric Born terms are of $\mathrm{O}(N_c)$ in the $1/N_c$ expansion, while those by the kaon electric Born terms are of $\mathrm{O}(N_c^0)$.

The Born terms through the magnetic coupling of the photon, which we call the magnetic Born terms, also contribute to the magnetic polarizabilities. The magnetic Born terms can interfere with the electric Born ones through the unitarity relation and contribute also to the magnetic polarizabilities. The magnetic Born term is written as the sum of the spin 1/2 and 3/2 baryon poles, each of which is of $O(N_c^{3/2})$. The large N_c consistency condition leads to the cancellation among the pole terms: The whole amplitude reduces to $O(N_c^{1/2})$, and as a result the amplitude is finite at infinite energies. We show that in order for the large N_c consistency condition to work in the pion production process off the Σ target the exotic strange baryon state denoted as Σ^{**} with the isospin 2 and spin 3/2, has to contribute to the amplitudes. Similarly, the condition requires two exotic states , Ξ_1^{**} and Ξ_3^{**} with isospin 3/2 and spin 1/2 and 3/2, respectively, for the kaon production amplitudes. The necessity of such exotic states is common to the large N_c baryon theories in order for the unitarity relations not to violate a definite N_c dependence of amplitudes, that is of $O(N_c^{1-n/2})$ for $n-$meson reaction amplitudes [19,20]. The contributions from the magnetic Born terms are of the same order as those from the electric ones in the $1/N_c$ expansion, but the magnetic Born contributions partially go beyond the calculation of $O(p^3)$ chiral order diagrams in HBChPT.

This paper is organized as follows: We discuss the pion and kaon electric Born contributions in the next section. The pion and kaon magnetic Born terms are given in Sec. III, and it is also discussed how the large N_c consistency condition works with the exotic states. The conclusions and discussion are given in the last section.

II. CONTRIBUTIONS FROM THE PION AND KAON ELECTRIC BORN TERMS

As stated in Introduction we use the dispersion integrals to compute the one-loop diagrams in the Compton scattering amplitudes, the imaginary parts of which are given through the electric Born terms.

A. Pion loop contributions

We start with the pion photoproduction amplitude, $T_a = \varepsilon_\mu T_a^\mu$ for $\gamma + Y \to \pi_a + Y'$ with $Y(Y')$ being the initial(final) hyperon with the strangeness $S = -1$, which is decomposed as

$$T_a = t_a^{(-)} T^{(-)} + t_a^{(+)} T^{(+)} + t_a^{(0)} T^{(0)} + t_a^{(\delta)} T^{(\delta)}, \tag{2.1}$$

where each amplitude is a function of the pion momentum $\boldsymbol{q}$ and photon $\boldsymbol{k}$, and the isospin factors are as follows: For $I = 1$ channel such as $\gamma\Sigma \to \pi\Sigma$, $t_a^{(-)} = i\varepsilon_{a3b} I_b$, $t_a^{(+)} = \{I_a, I_3\}$, $t_a^{(0)} = I_a$, $t_a^{(\delta)} = \delta_{a3} I$ with I_a being the conventional isospin matrix, and I the 3×3 unit matrix, and for $\gamma\Lambda \to \pi\Sigma$, $t_a^{(-)} = i\varepsilon_{a3b} \mathcal{T}_b^{\Sigma\Lambda}$ and $t_a^{(+)} = 0$ with $(\mathcal{T}_b^{\Sigma\Lambda})_{m0} = \delta_{mb}$.

The electric Born term of $\mathrm{O}(N_c^{1/2})$, that is of leading order in the $1/N_c$ expansion, is written model-independently as

$$T_E^{(-)} = \left(\frac{ef_{Y'Y\pi}}{4\pi m_\pi}\right)\left[i\boldsymbol{\sigma}\cdot\boldsymbol{\varepsilon} + 2i\boldsymbol{\sigma}\cdot\boldsymbol{t}\frac{\boldsymbol{\varepsilon}\cdot\boldsymbol{q}}{m_\pi^2-(q-k)^2}\right], \tag{2.2}$$

where $\boldsymbol{t} = \boldsymbol{k} - \boldsymbol{q}$, and $\boldsymbol{\varepsilon}$ is the polarization vector. Since other electric Born terms $T_E^{(+,0)}$ are of $\mathrm{O}(N_c^{-1/2})$ and break the unitarity limit at high energy, we ignore them hereafter [11].

The pseudovector pion coupling constants to the hyperons $f_{YY'\pi}$ are given as

$$\frac{1}{m_\pi}\frac{f_{Y'Y\pi}}{\sqrt{4\pi}} = \Lambda_{Y'Y}G_\pi \tag{2.3}$$

in the large N_c baryon model and in the BKSM [21], where the overall constant G_π is given empirically in the former and given in terms of the chiral angle $F(r)$ of the Skyrme model in the latter. The factor $\Lambda_{Y'Y}$ satisfies the following spin-flavor symmetry relation[1]:

$$\Lambda_{\Lambda\Sigma} = -\Lambda_{\Sigma\Sigma} = -\frac{1}{\sqrt{3}}\Lambda_{\Sigma^*\Lambda} = -\frac{2}{\sqrt{3}}\Lambda_{\Sigma^*\Sigma} = -\Lambda_{NN} = \frac{1}{3}. \tag{2.4}$$

This relation will play an important role for the large N_c consistency condition in the magnetic Born terms. To fix the pseudovector coupling constant we adopt $|f_{\Sigma\Lambda\pi}|/\sqrt{4\pi} = 0.22$ [21], which is close to the empirical value 0.20 ± 0.01 [22], and other coupling constants are obtained according to the Λ factor in the above.

According to Ref. [10–12] the forward dispersion relation with use of the electric Born term is known to give the electromagnetic polarizabilities as follows:

$$\begin{pmatrix} \alpha_Y(Y') \\ \\ \beta_Y^E(Y') \end{pmatrix} = A_{Y'}\left(\frac{ef_{Y'Y\pi}}{4\pi}\right)^2\frac{1}{24m_\pi^3}\begin{pmatrix} f(d) \\ \\ g(d) \end{pmatrix}, \tag{2.5}$$

where the factor $A_{Y'}$ is the multiplicity coming from the sum over a and the spin components. The functions $f(d)$ and $g(d)$ are defined as for $d > -1$ and $d \neq 1$ with $d = (M_{Y'} - M_Y)/m_\pi$

$$\begin{pmatrix} f(d) \\ \\ g(d) \end{pmatrix} = \frac{2}{\pi}\begin{pmatrix} d - 2A(d) + \dfrac{9(d+2A(d))}{(d^2-1)} \\ \\ -(d+2A(d)) \end{pmatrix}, \tag{2.6}$$

where

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[1] The sign of $\Lambda_{\Sigma\Sigma}$ and $\Lambda_{\Sigma^*\Sigma}$ is different from those in [21], because the sign of the isospin matrix for $I = 1$ is changed to the usual one, here.

$$A(d)=\begin{cases}\dfrac{1}{\sqrt{1-d^2}}\left[\tan^{-1}\dfrac{d}{\sqrt{1-d^2}}-\tan^{-1}\dfrac{1+d}{\sqrt{1-d^2}}\right] & (|d|<1)\\ -\dfrac{1}{2\sqrt{d^2-1}}\log\left(d+\sqrt{d^2-1}\right) & (d>1)\end{cases}, \tag{2.7}$$

and for $d=1$ we have $f(d)=16/\pi$, $g(d)=0$. In the above $\Sigma^*(1385)$ is also included in Y'. We use the empirical mass spectrum for the hyperons, the nucleon and Δ throughout the paper. The calculated results of the electromagnetic polarizabilities from the pion electric Born terms are tabulated in Table I.

We observe from Table I that the electric polarizability of the hyperons are in order

$$\alpha_{\Sigma^\pm} > \alpha_\Lambda > \alpha_{\Sigma^0}. \tag{2.8}$$

Due to the large coupling constant $f_{\Sigma^*\Lambda\pi}$, the contribution from Σ^* to the Λ electric polarizability is rather large, similar to the nucleon case. We note that the effect by the mass difference among the hyperons is rather significant as seen in the difference between the $\Lambda\pi$ contribution in the Σ target and $\Sigma\pi$ in the Λ target, where the former is exothermal and the latter endothermal. The difference between α_{Σ^+} and α_{Σ^-} cannot be calculated within this approximation.

B. Kaon loop contributions

We obtain the kaon electric Born term of $O(N_c^0)$ at leading order in the $1/N_c$ expansion for $\gamma+Y\to\bar{K}_\alpha(K_\alpha)+B$ with B being N and Δ (Ξ and Ξ^*) by replacing m_π by m_K and the coupling constant $f_{Y'Y\pi}$ by f_{YBK} in Eq.(2.2).

The P-wave kaon coupling constant f_{YBK} is of $O(N_c^0)$ and given as

$$\frac{1}{m_K}\frac{f_{YBK}}{\sqrt{4\pi}}=\Lambda_{YBK}G_K, \tag{2.9}$$

and we fix the pseudovector kaon coupling constant of ΛpK^- as $f_{\Lambda pK}/\sqrt{4\pi}=0.92$ [21], while the empirical one is 0.89 ± 0.10. The value 0.92 in the kaon mass scale as in Eq.(2.9) is reduced to 0.26 in the case of the pion mass scale, that is of the same order as the pion coupling constant. The large N_c relation of Λ_{YBK} for the charged kaons is given as

$$\Lambda_{\Lambda pK^-}=\sqrt{3}\Lambda_{\Lambda\Xi^-K^+}=-\frac{1}{2}\Lambda_{\Lambda\Xi^{*-}K^+}=-\frac{1}{\sqrt{2}} \tag{2.10}$$

for Λ vertices, and

$$\Lambda_{\Sigma^-nK^-}=-\frac{1}{\sqrt{3}}\Lambda_{\Sigma^-\Delta^0K^-}=-\frac{\sqrt{3}}{5}\Lambda_{\Sigma^+\Xi^0K^+}=-\frac{1}{2}\Lambda_{\Sigma^+\Xi^{*0}K^+}=-\frac{1}{3} \tag{2.11}$$

for Σ vertices.

The kaon contributions are given as

$$\begin{pmatrix} \alpha_Y \\ \\ \beta_Y^E \end{pmatrix} = \sum_B A_B \left(\frac{e f_{YBK}}{4\pi}\right)^2 \frac{1}{24 m_K^3} \begin{pmatrix} f(d) \\ \\ g(d) \end{pmatrix}, \tag{2.12}$$

where A_B is the same multiplicity as $A_{Y'}$ in the pion production, and $d = (M_Y - M_B)/m_K$. Note that the factor $(m_\pi/m_K) = 0.279$ reduces the size of the kaon contributions. The numerical results from the kaon electric Born terms are tabulated in Table II.

The kaon-loop contributions to polarizabilities are in order

$$\alpha_\Lambda > \alpha_{\Sigma^+} > \alpha_{\Sigma^0} > \alpha_{\Sigma^-}. \tag{2.13}$$

We see that the contributions from the decuplet baryons in the final states are of the same order as those from the octet baryons.

The kaon contribution leads to the result $\alpha_{\Sigma^+} > \alpha_{\Sigma^-}$, that is of the same sign in Ref. [14]. At the same time the kaon contribution to the nucleon makes the proton polarizabilities larger than the neutron ones, but it does not agree with the experimental tendency: $\alpha_n > \alpha_p$.

III. MAGNETIC BORN TERMS AND THE LARGE N_C CONSISTENCY CONDITION

The spatial part of the electromagnetic current, $\boldsymbol{J}$, contributes to the magnetic Born term, where

$$< Y'(\boldsymbol{p}')|\boldsymbol{\varepsilon}\cdot\boldsymbol{J}|Y(\boldsymbol{p})> = < Y'|i\boldsymbol{s}\cdot\boldsymbol{\mu}|Y> \tag{3.1}$$

with $\boldsymbol{s} = (\boldsymbol{p}' - \boldsymbol{p}) \times \boldsymbol{\varepsilon}$ and $\boldsymbol{\mu}$ being the magnetic moment operator. The magnetic moment is decomposed as

$$< Y'|\boldsymbol{\mu}|Y> = \boldsymbol{S}\frac{e}{2M_N}(\mu^V_{Y'Y}\mathcal{T}_3 + \mu^S_{Y'Y}), \tag{3.2}$$

where $\boldsymbol{S}(\mathcal{T}_3)$ is the transition spin (isospin) matrix, and $\mu^V_{Y'Y}$ ($\mu^S_{Y'Y}$) is the isovector (isoscalar) part of the hyperon magnetic moment in units of Bohr magneton. Since the isovector part, $\mu^V e/2M_N$, is of O(N_c), while the isoscalar part is of O(N_c^0), the leading contributions come from the isovector part of the magnetic moments. We further note that $\mu^V_{Y'Y}$ having the same strangeness S is proportional to the factor $\Lambda_{Y'Y}$ of Eq.(2.4) [19,23,24].

The magnetic Born term for a process $\gamma + Y \to \pi^a + Y'$ is written as

$$T_M^a = \sum_{Y''} \left\{ \left(\frac{e f_{Y'Y''\pi}}{8\pi m_\pi M_N} \right) \mu^V_{Y''Y} \frac{(\boldsymbol{S}' \cdot \boldsymbol{q})^\dagger (\boldsymbol{S} \cdot \boldsymbol{s}) \mathcal{T}_a'^\dagger \mathcal{T}_3}{M_{Y''} - M_{Y'} - \omega_q} + \left(\frac{e f_{Y''Y\pi}}{8\pi m_\pi M_N} \right) \mu^V_{Y'Y''} \frac{(\boldsymbol{S}' \cdot \boldsymbol{s})^\dagger (\boldsymbol{S} \cdot \boldsymbol{q}) \mathcal{T}_3'^\dagger \mathcal{T}_a}{M_{Y''} - M_Y + \omega_q} \right\} \tag{3.3}$$

at leading order of the $\mathrm{O}(1/N_c)$ expansion, where $\boldsymbol{S}'(\mathcal{T}_a')$ is the transition spin(isospin) matrix for the $Y'' \to Y'$ vertex, while those without a prime for the $Y \to Y''$ vertex. The magnetic Born terms are decomposed as

$$T_M^a = \sum_{\ell=1,3} \sum_{n=\pm,\delta} \mathcal{P}_\ell(\hat{\boldsymbol{q}}, \hat{\boldsymbol{s}})\, t_a^{(n)}\, T_\ell^{(n)}(\omega_q), \tag{3.4}$$

where $\mathcal{P}_1(\mathcal{P}_3)$ is the projection operator for the P-wave pion production amplitude with total angular momentum $j = 1/2(3/2)$:

$$\mathcal{P}_1(\hat{\boldsymbol{q}}, \hat{\boldsymbol{s}}) = (\boldsymbol{\sigma} \cdot \hat{\boldsymbol{q}})(\boldsymbol{\sigma} \cdot \hat{\boldsymbol{s}}), \quad \text{and} \quad \mathcal{P}_3(\hat{\boldsymbol{q}}, \hat{\boldsymbol{s}}) = 3(\hat{\boldsymbol{q}} \cdot \hat{\boldsymbol{s}}) - (\boldsymbol{\sigma} \cdot \hat{\boldsymbol{q}})(\boldsymbol{\sigma} \cdot \hat{\boldsymbol{s}}) \tag{3.5}$$

with $\hat{\boldsymbol{q}} = \boldsymbol{q}/q$ and $\hat{\boldsymbol{s}} = \boldsymbol{s}/k$ for Y' with spin 1/2. Similar expressions are written for Y' with spin 3/2. The final states are restricted to the states with an octet or decuplet baryon accompanied with a pion or kaon.

We notice here that in order for the large N_c consistency condition to hold the exotic hyperon states are required for the Σ target. Due to the consistency condition the magnetic Born terms keep being of $\mathrm{O}(N_c^{1/2})$ and are convergent at infinite energies, as a result. So, we concentrate ourselves to the Σ target both for the pion and kaon magnetic Born terms in the following.

A. The pion magnetic Born terms

Here we discuss explicitly the magnetic Born terms of the process $\gamma + \Sigma \to \pi^a + \Sigma$. Using the mass abbreviation $\Delta_{Y'Y} = M_{Y'} - M_Y$, we write the non-exotic pole amplitudes with the Σ and Σ^* poles in this channel as

$$T_1^{(\pm)} = \left(\frac{e f_{\Sigma\Sigma\pi}}{8\pi m_\pi M_N} \right) \mu^V_\Sigma \left[-\frac{kq}{2\omega_q} \mp \frac{kq}{6\omega_q} \right] \pm \frac{1}{3} \left(\frac{e f_{\Sigma^*\Sigma\pi}}{8\pi m_\pi M_N} \right) \mu^V_{\Sigma^*\Sigma} \left[\frac{2}{3} \frac{kq}{\Delta_{\Sigma^*\Sigma} + \omega_q} \right], \tag{3.6}$$

$$T_3^{(\pm)} = \left(\frac{e f_{\Sigma\Sigma\pi}}{8\pi m_\pi M_N} \right) \mu^V_\Sigma \left[\pm \frac{kq}{3\omega_q} \right] + \frac{1}{3} \left(\frac{e f_{\Sigma^*\Sigma\pi}}{8\pi m_\pi M_N} \right) \mu^V_{\Sigma^*\Sigma} \left[\frac{1}{2} \frac{kq}{\Delta_{\Sigma^*\Sigma} - \omega_q} \pm \frac{1}{6} \frac{kq}{\Delta_{\Sigma^*\Sigma} + \omega_q} \right], \tag{3.7}$$

and those with the Λ pole are similarly written, but not shown explicitly. We note that each pole term is of $\mathrm{O}(N_c^{3/2})$.

If we use the relation given by the Λ factors,

$$f_{\Sigma\Sigma\pi}\mu^V_{\Sigma} = \frac{4}{3}f_{\Sigma^*\Sigma\pi}\mu^V_{\Sigma^*\Sigma} = f_{\Sigma\Lambda\pi}\mu^V_{\Sigma\Lambda}, \tag{3.8}$$

we see that the cancellation does not occur for the sums of the above amplitudes. Notice, however, that the $\Sigma+\pi$ channel can communicate with the $I=2$ channel with strangeness -1. The large N_c baryon theories and BKSM predict two such exotic baryons with spin 3/2 and 5/2, where the P-wave antikaon is bound around the soliton with isospin 2 in the latter model. In this channel we need the exotic state with $I=2$ and spin 3/2 in order for the large N_c consistency condition to hold. (The exotic state with spin 5/2 cannot interact with P-wave $\pi\Sigma$ states.) We denote the exotic state with spin 3/2 as Σ^{**}. Including the exotic state, we can see that the cancellation occurs with the spin-flavor symmetry relation

$$f_{\Sigma^{**}\Sigma\pi}\mu^V_{\Sigma^{**}\Sigma} = \frac{3}{2}f_{\Sigma^*\Lambda\pi}\mu^V_{\Sigma^*\Lambda} = \frac{9}{2}f_{\Sigma\Sigma\pi}\mu^V_{\Sigma} = \frac{9}{2}f_{\Sigma\Lambda\pi}\mu^V_{\Sigma\Lambda}. \tag{3.9}$$

Indeed, BKSM gives the factor $\Lambda_{\Sigma^{**}\Sigma}$ as

$$\Lambda_{\Sigma^{**}\Sigma} = -\frac{1}{\sqrt{2}}, \tag{3.10}$$

that is consistent with the above condition for the cancellation, of course. This result is shared with the large N_c baryon theories.

We summarize the resultant Born terms as

$$T_1^{(-)} = \left(\frac{ef_{\Lambda\Sigma\pi}}{8\pi m_\pi M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[-\frac{1}{6}\frac{\Delta_{\Sigma^*\Lambda}}{(\Delta_{\Sigma^*\Sigma}+\omega_q)(\Delta_{\Sigma\Lambda}-\omega_q)} - \frac{1}{3}\frac{\Delta_{\Sigma^{**}\Sigma}}{\omega_q(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right.$$
$$\left. - \frac{1}{2}\frac{(\Delta_{\Sigma^{**}\Sigma}-\Delta_{\Sigma\Lambda})}{(\Delta_{\Sigma\Lambda}+\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right], \tag{3.11}$$

$$T_1^{(+)} = \left(\frac{ef_{\Lambda\Sigma\pi}}{8\pi m_\pi M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[-\frac{1}{2}\frac{\Delta_{\Sigma\Lambda}}{\omega_q(\Delta_{\Sigma\Lambda}+\omega_q)} - \frac{1}{6}\frac{\Delta_{\Sigma^*\Sigma}}{\omega_q(\Delta_{\Sigma^*\Sigma}+\omega_q)}\right.$$
$$\left. + \frac{1}{6}\frac{\Delta_{\Sigma^{**}\Lambda}}{(-\Delta_{\Sigma\Lambda}+\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right], \tag{3.12}$$

$$T_1^{(\delta)} = \left(\frac{ef_{\Lambda\Sigma\pi}}{8\pi m_\pi M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[-\frac{\Delta_{\Sigma^{**}\Sigma}-\Delta_{\Sigma\Lambda}}{(\Delta_{\Sigma\Lambda}+\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)} + \frac{1}{3}\frac{\Delta_{\Sigma^{**}\Lambda}}{(\Delta_{\Sigma\Lambda}-\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right], \tag{3.13}$$

$$T_3^{(-)} = \left(\frac{ef_{\Lambda\Sigma\pi}}{8\pi m_\pi M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[\frac{1}{4}\frac{\Delta_{\Sigma^*\Sigma}}{({\Delta_{\Sigma^*\Sigma}}^2-\omega_q^2-i\Delta_{\Sigma^*\Sigma}\Gamma_{\Sigma^*})} - \frac{5}{4}\frac{\Delta_{\Sigma^{**}\Sigma}}{({\Delta_{\Sigma^{**}\Sigma}}^2-\omega_q^2-i\Delta_{\Sigma^{**}\Sigma}\Gamma_{\Sigma^{**}})}\right.$$
$$\left. - \frac{1}{3}\frac{\Delta_{\Sigma^{**}\Sigma}}{\omega_q(\Delta_{\Sigma^{**}\Sigma}+\omega_q)} - \frac{1}{6}\frac{\Delta_{\Sigma^{**}\Sigma^*}}{(\Delta_{\Sigma^*\Sigma}+\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)} - \frac{1}{3}\frac{\Delta_{\Sigma^{**}\Lambda}}{(-\Delta_{\Sigma\Lambda}+\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right], \tag{3.14}$$

$$T_3^{(+)} = \left(\frac{ef_{\Lambda\Sigma\pi}}{8\pi m_\pi M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[\frac{1}{4}\frac{\Delta_{\Sigma^*\Sigma}}{({\Delta_{\Sigma^*\Sigma}}^2-\omega_q^2-i\Delta_{\Sigma^*\Sigma}\Gamma_{\Sigma^*})} - \frac{1}{4}\frac{\Delta_{\Sigma^{**}\Sigma}}{({\Delta_{\Sigma^{**}\Sigma}}^2-\omega_q^2-i\Delta_{\Sigma^{**}\Sigma}\Gamma_{\Sigma^{**}})}\right.$$
$$\left. + \frac{1}{3}\frac{\Delta_{\Sigma\Lambda}}{\omega_q(\Delta_{\Sigma\Lambda}-\omega_q)} - \frac{1}{12}\frac{\Delta_{\Sigma^{**}\Sigma^*}}{(\Delta_{\Sigma^*\Sigma}+\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right], \tag{3.15}$$

$$T_3^{(\delta)} = \left(\frac{ef_{\Lambda\Sigma\pi}}{8\pi m_\pi M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[\frac{2\Delta_{\Sigma^{**}\Sigma}}{{\Delta_{\Sigma^{**}\Sigma}}^2-\omega_q^2-i\Delta_{\Sigma^{**}\Sigma}\cdot\Gamma_{\Sigma^{**}}} - \frac{2}{3}\frac{\Delta_{\Sigma^{**}\Lambda}}{(\Delta_{\Sigma\Lambda}-\omega_q)(\Delta_{\Sigma^{**}\Sigma}+\omega_q)}\right]. \tag{3.16}$$

Since the mass differences of $\mathrm{O}(N_c^{-1})$ appear in the numerators, the resultant amplitudes are reduced to $\mathrm{O}(N_c^{1/2})$ and finite at high energies. We stress here that we have not introduced any vertex functions depending on the meson or photon momentum, because the vertex corrections go beyond the one-loop approximation. In this sense the exotic states play a role similar to a natural cutoff function without destroying the analytic structure of the one-loop amplitudes.

In the above we inserted the total width Γ_{Σ^*} and $\Gamma_{\Sigma^{**}}$ into the pure resonance terms of T_3. The width Γ_{Σ^*} is given as

$$\Gamma_{\Sigma^*} = \frac{2}{3}\left(\frac{f^2_{\Sigma^*\Lambda\pi}}{4\pi}\right)\frac{q^3_\Lambda}{m^2_\pi} + \frac{4}{3}\left(\frac{f^2_{\Sigma^*\Sigma\pi}}{4\pi}\right)\frac{q^3_\Sigma}{m^2_\pi}, \tag{3.17}$$

where $q_\Lambda(q_\Sigma)$ is the pion momentum decaying into the channel $\Lambda(\Sigma) + \pi$. This form of the width is the same as that adopted in the previous works [9,12], which guarantees the narrow width limit. Numerical values for the widths including recoil are $\Gamma^{\Lambda\pi}_{\Sigma^*} = 44.8$ MeV and $\Gamma^{\Sigma\pi}_{\Sigma^*} = 4.7$ MeV. Similar form of the width of Σ^{**} is used, where Σ^{**} is supposed to decay only to $\Sigma\pi$ channel, because the Σ^{**} mass is not expected to be so higher than the $\Sigma^*\pi$ threshold as seen below. Then, we have $\Gamma^{\Sigma\pi}_{\Sigma^{**}} = 138$ MeV, but since $\Lambda + 2\pi$ channel is not taken into account, though it opens, the width of Σ^{**} would be underestimated.

Since all the isovector magnetic moments appearing in the amplitudes are for the hyperons with $S = -1$, we take the empirical value, $\mu^V_{\Sigma\Lambda} = -1.61$, to fix the magnetic coupling constants. The isoscalar magnetic moment can give pole terms of $\mathrm{O}(N_c^{1/2})$, but they are not of leading order. Since the cancellation among the pole terms does not hold at non-leading order and then the unitarity bound is broken, we disregard them as in the case of the electric pion Born terms.

B. The kaon magnetic Born terms

Contrary to the pion photoproduction processes, the isovector magnetic moments with different strangeness from $S = 0$ to -2, contribute to the $\bar{K}$ or K photoproduction amplitudes. Although the magnetic moment is not completely proportional to the Λ factor, we take the experimental value for $\mu^V_{\Sigma\Lambda}$, by which the other magnetic moments are fixed, because it may be regarded as giving an average.

In the case of the $\gamma + \Sigma \rightarrow K + \Xi$ process, the cancellation does not occur among the Ξ and Ξ^* pole terms: The two exotic Ξ states with isospin 3/2 contribute to this processes, the one with spin 1/2 denoted as Ξ^{**}_1 and the other with spin 3/2 as Ξ^{**}_3. The resultant amplitudes are written as

$$T_1 = \left(\frac{ef_{\Sigma^+\Xi^0K^+}}{8\pi m_K M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[\frac{16}{45}\frac{\Delta_{\Sigma^*\Xi}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Xi^*\Sigma}+\omega_q)}\right.$$
$$\left.+\ \frac{4}{45}\frac{\Delta_{\Xi_1^{**}\Xi}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Xi_1^{**}\Sigma}+\omega_q)}+\frac{4}{9}\frac{\Delta_{\Xi_3^{**}\Xi}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Xi_3^{**}\Sigma}+\omega_q)}\right], \tag{3.18}$$
$$T_3 = \left(\frac{ef_{\Sigma^+\Xi^0K^+}}{8\pi m_K M_N}\right)\mu^V_{\Sigma\Lambda}\cdot kq\left[\frac{1}{5}\frac{\Delta_{\Sigma^*\Sigma}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Sigma^*\Xi}-\omega_q)}\right.$$
$$\left.+\ \frac{4}{45}\frac{\Delta_{\Sigma^*\Xi}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Xi^*\Sigma}+\omega_q)}-\frac{8}{45}\frac{\Delta_{\Xi_1^{**}\Xi}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Xi_1^{**}\Sigma}+\omega_q)}+\frac{1}{9}\frac{\Delta_{\Xi_3^{**}\Xi}}{(\Delta_{\Xi\Sigma}+\omega_q)(\Delta_{\Xi_3^{**}\Sigma}+\omega_q)}\right] \tag{3.19}$$

for $\gamma + \Sigma^+ \rightarrow K^+ + \Xi^0$, and similar amplitudes are written for other Σ targets. In the above we used

$$\Lambda_{\Sigma^+\Xi_1^{**0}K^+} = \frac{\sqrt{2}}{3\sqrt{3}} \quad \text{and} \quad \Lambda_{\Sigma^+\Xi_3^{**0}K^+} = \frac{\sqrt{5}}{3}, \tag{3.20}$$

which are predicted by BKSM. Each pole term in the above is of $\mathrm{O}(N_c)$, but the resultant amplitudes are of $\mathrm{O}(N_c^0)$, because the mass differences appearing in the numerators are of $\mathrm{O}(N_c^{-1})$. The large N_c consistency condition also works to reduce the N_c dependence and to converge the asymptotic behavior. Since the resonance poles do not develop in the physical region of ω_q, their contributions are rather small.

C. Mass spectrum of the exotic states

The exotic states are required to satisfy the large N_c consistency condition of the production amplitudes as shown in previous subsections. Here we estimate the masses within BKSM. The baryon mass spectrum of the model is given by the following formula as

$$M = M_s + |s|\omega + \frac{1}{2\Lambda}\left[cJ(J+1) + (1-c)I(I+1) - c(1-c)\frac{|s|}{2}\left(\frac{|s|}{2}+1\right)\right], \tag{3.21}$$

where M_s, ω, c and Λ are the parameters to be calculated by the model [18,23]. Instead of calculating these parameters by the model we estimate them by the existing mass spectrum of the non-strange and strange baryons: The result we adopt is $M_s = 866$ MeV, $c = 0.630$, $\omega = 221$ MeV and $M_\Delta - M_N = 3/(2\Lambda) = 293$ MeV. The same parameters give the masses of the exotic states as

$$\begin{aligned}
\Sigma^{**}(I=2, J=\tfrac{3}{2}) &= 1517\,\mathrm{MeV}, \\
\Xi_1^{**}(I=\tfrac{3}{2}, J=\tfrac{1}{2}) &= 1444\,\mathrm{MeV}, \qquad \Xi_3^{**}(I=\tfrac{3}{2}, J=\tfrac{3}{2}) = 1639\,\mathrm{MeV}.
\end{aligned} \tag{3.22}$$

We note that almost the same values are obtained for the masses of the exotic states by the mass formula in the tree level of the large N_c chiral perturbation theory [25].

The mass of Σ^{**} is high enough to decay into $\Sigma\pi, \Sigma\pi\pi$ and $\Lambda\pi\pi$ channels, while the mass of Ξ_1^{**} is low and seems to be stable. It should be noticed, however, that the exotic states could disappear at $N_c = 3$. The masses of the exotic states may be sensitive to higher N_c corrections, but we use 1520 MeV for Σ^{**}, 1450 (1640) MeV for $\Xi_1^{**}(\Xi_3^{**})$ in this paper.

D. Magnetic polarizability from the magnetic Born terms

The magnetic polarizability β_Y^M is given by the integration over energy as follows:

$$\beta_Y^M = \frac{2}{\pi}\int_{\omega_{\rm th}}^{\infty}\frac{d\omega_k}{\omega_k^2}\frac{q}{\omega_k}\sum_{m,n}\left(T_1^{(n)^*}T_1^{(m)} + 2T_3^{(n)^*}T_3^{(m)}\right)\sum_a t_a^{(n)^\dagger}t_a^{(m)} \tag{3.23}$$

for the spin 1/2 final baryon.

The magnetic Born term can interfere with the electric one: By the angular integration we have

$$\beta_Y^I = \sum_{Y'}\frac{2G_E}{\pi}\int\frac{d\omega_k}{\omega_k^2}\frac{q}{\omega_k}\left[\frac{1}{v} - \frac{1-v^2}{2v^2}\log\left(\frac{1+v}{1-v}\right)\right]\sum_{n,a}\left({\rm Re}T_1^{(n)} - {\rm Re}T_3^{(n)}\right)t_a^{(-)^\dagger}t_a^{(n)} \tag{3.24}$$

for the spin 1/2 final baryon, where G_E denotes the corresponding coupling constant in T_E. In the case of the spin 3/2 final baryon, the above expressions are little changed. In Table III we show the numerical results of the magnetic polarizabilities, in which all the contributions are included.

Instead of integrating the full amplitudes, if we pick up only the the Σ^* state and ignore the exotic state and the background contributions at all, we may get rid of the contributions from the exotic state. Such a narrow width approximation has been discussed in the case of the nucleon polarizabilities and shown to give the same result as the one by HBChPT in Ref. [9,12]. So, we proceed to the narrow width approximation for the $\gamma + \Lambda \to \Lambda + \pi$ channel as a typical example: $|T_3^{(\pm)}|^2$ contains the Σ^* resonance and its contribution in Eq.(3.23) is proportional to

$$\frac{e^2}{4\pi}\left(\frac{\mu_{\Sigma\Lambda}^V}{2M_N}\right)^2\frac{2}{\pi}\int\frac{d\omega_k}{\omega_k^3}\left(\frac{f_{\Sigma\Lambda\pi}^2}{4\pi}\frac{q^3}{m_\pi^2}\right)\frac{4{\Delta_{\Sigma^*\Lambda}}^2}{({\Delta_{\Sigma^*\Lambda}}^2 - \omega_q^2)^2 + (\Delta_{\Sigma^*\Lambda}\,\Gamma_{\rm tot})^2}, \tag{3.25}$$

where $\Gamma_{\rm tot}$ is the total width of Σ^*. Using the relation

$$\frac{f_{\Sigma\Lambda\pi}^2}{4\pi}\frac{q^3}{m_\pi^2} = \frac{1}{2}\left[\frac{2}{3}\frac{f_{\Sigma^*\Lambda\pi}^2}{4\pi}\frac{q^3}{m_\pi^2}\right] = \frac{1}{2}\Gamma_{\Lambda\pi}, \tag{3.26}$$

and taking the narrow width limit as

$$\lim_{\Gamma_{\rm tot}\to 0}\frac{\Delta_{\Sigma^*\Lambda}\Gamma_{\rm tot}}{({\Delta_{\Sigma^*\Lambda}}^2 - \omega_q^2)^2 + (\Delta_{\Sigma^*\Lambda}\,\Gamma_{\rm tot})^2} = \pi\delta({\Delta_{\Sigma^*\Lambda}}^2 - \omega_q^2), \tag{3.27}$$

we may have

$$\beta^M_{\Lambda\to\Lambda}|_{\Sigma^*} = \frac{e^2}{4\pi}\left(\frac{\mu^V_{\Sigma\Lambda}}{2M_N}\right)^2 \frac{4}{\Delta_{\Sigma^*\Lambda}}\left(\frac{\Gamma_{\Lambda\pi}}{\Gamma_{\rm tot}}\right), \tag{3.28}$$

where the spin factor 2 is multiplied and the last factor is the branching ratio of $\Sigma^* \to \Lambda\pi$. Adding the $\Sigma^\pm\pi^\mp$ channels, we get

$$\beta^M_\Lambda|_{\Sigma^*} = \left(\frac{e^2}{4\pi}\right)\left(\frac{\mu^V_{\Sigma\Lambda}}{2M_N}\right)^2 \frac{4}{\Delta_{\Sigma^*\Lambda}}, \tag{3.29}$$

which is similar to the case of the nucleon. For the Σ target we have

$$\beta^M_{\Sigma^\pm}|_{\Sigma^*} = \left(\frac{e^2}{4\pi}\right)\left(\frac{\mu^V_{\Sigma\Lambda}}{2M_N}\right)^2 \frac{1}{\Delta_{\Sigma^*\Sigma}}, \qquad \beta^M_{\Sigma^0}|_{\Sigma^*} = 0. \tag{3.30}$$

Since Σ^0 does not have the leading isovector magnetic moment, $\beta_{\Sigma^0}|_{\rm narrow}$ is zero at leading order. Even if the isoscalar magnetic moment is used, it is at most 1/4 of the $\beta_{\Sigma^\pm}$, because $\mu^S \approx 1/2\mu^V_\Sigma$. The numerical results are as follows:

$$\beta^M_\Lambda|_{\Sigma^*} = 6.13 \quad \text{and} \quad \beta^M_{\Sigma^\pm}|_{\Sigma^*} = 2.15$$

in units of $10^{-4}\,\mathrm{fm}^3$.

Similar narrow width approximation to Σ^{**} gives the values,

$$\beta^M_{\Sigma^\pm}|_{\Sigma^{**}} = 3.78 \quad \text{and} \quad \beta^M_{\Sigma^0}|_{\Sigma^{**}} = 5.04.$$

Since Γ^{**}_Σ is broad as seen previously, these values would be an overestimate, but the exotic resonance contributions cannot be discarded, especially to Σ^0.

Finally, we note that β in the narrow width approximation is of $\mathrm{O}(N_c^3)$, because the limit picks up only the relevant pole of $\mathrm{O}(N_c^{3/2})$; that is, the limit is not consistent with the $1/N_c$ expansion.

IV. CONCLUSIONS AND DISCUSSION

We have calculated the spin-averaged electromagnetic polarizabilities of the hyperons Λ and Σ within the one-loop approximation. In order to calculate the one-loop diagrams we used the dispersion relations, where the imaginary parts are given by the Born terms in the pion and kaon photoproduction Born terms. The Born terms satisfy the low energy theorems, and their form is model-independent. The coupling constants are determined so as to satisfy spin-flavor symmetry of the large N_c QCD.

The calculated electromagnetic polarizabilities through the pion and kaon Born terms are summarized as $\alpha_\Lambda = 18.05$, $\alpha_{\Sigma^+} = 22.08$, $\alpha_{\Sigma^0} = 13.79$ and $\alpha_{\Sigma^-} = 18.71$, and $\beta_\Lambda = 3.22$, $\beta_{\Sigma^+} = 6.67$, $\beta_{\Sigma^0} = 5.52$ and $\beta_{\Sigma^-} = 7.13$ in units of $10^{-4}\,\mathrm{fm}^3$. These values would be too large as seen from the large values of the polarizabilities of the nucleon given by the same calculation [12] as well as the one-loop calculation in HBChPT [3]. This is because the high energy contributions from the one-loop diagrams are not fully reduced for the spin-averaged polarizabilities compared to the spin polarizabilities owing to the power behavior of the energy denominator in the integrals. In order to reduce the values of the spin-averaged polarizabilities within the one-loop calculations we would have to go to the approximation beyond the one-loop level, such as vertex corrections and the unitarization of the Born amplitudes.

The electric Born terms would give the same spin-averaged polarizabilities as the SU(3) extension of HBChPT [14], if the hyperon mass differences are ignored. But we observed that the polarizabilities strongly depend on the hyperon mass difference, and then SU(3) symmetry of the polarizabilities is further broken besides the symmetry breaking due to the pion and kaon mass difference, even if the coupling constants satisfy SU(3) symmetry with an appropriate F/D ratio.

As to the magnetic Born amplitudes we have shown that exotic hyperon states are inevitably required even in the non-exotic reaction channel in order for the large N_c consistency condition to hold. The consistency condition guarantees meson-baryon reaction amplitudes to have a meaningful N_c limit. Due to the consistency condition the magnetic Born terms remain at $\mathrm{O}(N_c^{1/2})$, and as a result they become finite at high energies and give finite magnetic polarizabilities as the electric Born ones. We also noted that the narrow width limit is not consistent with the $1/N_c$ expansion, because the limit picks up the single resonance pole term of $\mathrm{O}(N_c^{3/2})$ and the resultant polarizability is of $\mathrm{O}(N_c^3)$. If we reduce N_c to 3, the coupling constants of the exotic hyperons to non-exotic ones would vanish, but simultaneously it makes the magnetic Born amplitudes break the unitarity bound at high energies even for such a case of the nucleon. Thus, it is impossible to have finite results within the one-loop approximation. Contrary, it is a serious problem for the large N_c baryon theories to study the $1/N_c$ corrections to physical quantities related to the exotic states and what physical effects are expected by the exotic states besides a role of the natural cutoff, if the leading terms in the $1/N_c$ expansion are valid. These tasks are left to further investigations.

It is known that there is a negative parity hyperon $\Lambda^*(1405)$, which BKSM predicts as an S-wave bound state of the antikaon around the $I = J = 0$ chiral SU(2) soliton. The model also predicts $\Sigma_{1/2^-}$ and $\Sigma_{3/2^-}$, which are the bound states to the $I = J = 1$ soliton [21]. The electric dipole transition amplitudes with the poles at the negative parity hyperons of spin 1/2 give the electric polarizabilities. Since the transition electric dipole moment is of $\mathrm{O}(N_c^0)$ and the S-wave pion coupling constant of $\mathrm{O}(N_c^{-1/2})$, the electric polarizabilities are of $\mathrm{O}(N_c^{-1})$, compared to the electric Born contributions of $\mathrm{O}(N_c)$. Taking the narrow width approximation given by

$$\alpha^D = \frac{e^2}{4\pi}\left(\frac{\kappa_{\Lambda^* Y}}{2M_N}\right)^2 \frac{2}{M_{\Lambda^*} - M_Y},$$

we get $\alpha_\Lambda^D = 0.18$ and $\alpha_{\Sigma^0}^D = 0.24$ in units of $10^{-4}\,\mathrm{fm}^3$, where we used the transition dipole moments $\kappa_{\Lambda^*\Lambda} = \kappa_{\Lambda^*\Sigma} = 0.41$ in units of the Bohr magneton, which are given by BKSM. The model also predicts that the dipole moment of $\Sigma^*_{1/2^-}$ is $-1/3$ of the Λ^*, and then the contribution to α would be tiny; $\alpha_{\Sigma^+}^D = 0.08$, $\alpha_{\Sigma^0}^D = 0.02$ and $\alpha_{\Sigma^-}^D = 0$. Our values are quite different from those of Ref. [13], but of the same order as Ref. [15]. The interference terms between the electric Born and the electric dipole moment terms are also small.

Gobbi et al. calculated the polarizabilities of the hyperons in BKSM [15], but they used the two-photon seagull terms in the Lagrangian. It is, however, pointed out that it is dangerous to use the two-photon-seagull terms in the Lagrangian to calculate the polarizabilities, because the gauge invariance makes the seagull terms vanish for the electric polarizability [11,26,27]. Although we referred to the same BKSM, our approach to the subject is quite different from theirs, and the results are also different: We point out that the chiral soliton model including BKSM gives the model-independent form of the pion and kaon photoproduction amplitudes at tree level and then the polarizabilities are given by calculating the loop integrals with the dispersion relations.

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TABLES

TABLE I. The electromagnetic polarizabilities from the pion electric Born terms in units of $10^{-4}\,\mathrm{fm}^3$. $f_{\Sigma\Lambda\pi}/\sqrt{4\pi}$ =0.22 is used to fix the pion coupling constants.

Y	Λ			$\Sigma^\pm$				Σ^0		
Y'	Σ	Σ^*	total	Λ	Σ	Σ^*	total	Σ	Σ^*	total
α_Y	5.40	7.13	12.54	11.89	4.30	0.98	17.17	8.61	1.95	10.56
β_Y^E	0.34	−1.29	−0.95	0.87	0.43	−0.07	1.23	0.86	−0.14	0.72

TABLE II. The electromagnetic polarizabilities from the kaon electric Born terms in units of $10^{-4}\,\mathrm{fm}^3$. $f_{\Lambda pK^-}/\sqrt{4\pi}$ =0.92 is used to fix the kaon coupling constants.

Y	Λ			Σ^+			Σ^0			Σ^-		
B	$N+\Xi$	Ξ^*	total	Ξ	$\Delta+\Xi^*$	total	$N+\Xi$	$\Delta+\Xi^*$	total	N	Δ	total
α_Y	3.16	2.36	5.52	2.36	2.56	4.91	1.62	1.61	3.23	0.88	0.67	1.55
β_Y^E	0.28	0.06	0.34	0.21	0.22	0.44	0.14	0.14	0.28	0.07	0.07	0.13

TABLE III. Total magnetic polarizability β and each contribution β_Y^E, β_Y^M or β_Y^I in units of $10^{-4}\,\mathrm{fm}^3$. $\mu_{\Sigma\Lambda}^V$=−1.61 is used to fix the isovector part of hyperon magnetic moments.

Y	total β	π or K-loop	β^E	β^M	β^I	sum
Λ	3.22	$\pi+\Lambda,\Sigma$	0.34	5.36	−0.73	4.97
		$\pi+\Sigma^*$	−1.29	0.24	−1.25	−2.31
		$\bar{K}+N, K+\Xi$	0.28	1.04	−0.58	0.74
		$K+\Xi^*$	0.06	0.07	−0.31	−0.18
Σ^+	6.67	$\pi+\Lambda,\Sigma$	1.30	5.08	0.24	6.61
		$\pi+\Sigma^*$	−0.07	0.13	−0.20	−0.13
		$\bar{K}+N, K+\Xi$	0.21	0.40	−0.34	0.27
		$\bar{K}+\Delta, K+\Xi^*$	0.22	0.14	−0.45	−0.08
Σ^0	5.52	$\pi+\Sigma$	0.86	4.06	0.46	5.38
		$\pi+\Sigma^*$	−0.14	0.52	−0.42	−0.03
		$\bar{K}+N, K+\Xi$	0.14	0.18	−0.19	0.13
		$\bar{K}+\Delta, K+\Xi^*$	0.14	0.16	−0.25	0.05
Σ^-	7.13	$\pi+\Lambda,\Sigma$	1.30	5.08	0.24	6.61
		$\pi+\Sigma^*$	−0.07	0.13	−0.20	−0.13
		$\bar{K}+N, K+\Xi$	0.07	0.40	0.04	0.51
		$\bar{K}+\Delta, K+\Xi^*$	0.07	0.14	−0.07	0.14